

May 2010

Stable mass hierarchies and dark matter from hidden sectors in the scale-invariant standard model

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Abstract

Scale invariance may be a classical symmetry which is broken radiatively. This provides a simple way to stabilise the scale of electroweak symmetry breaking against radiative corrections. But for such a theory to be fully realistic, it must actually incorporate a hierarchy of scales, including the Planck and the neutrino mass scales in addition to the electroweak scale. The dark matter sector and the physics responsible for baryogenesis may or may not require new scales, depending on the scenario. We develop a generic way of using hidden sectors to construct a technically-natural hierarchy of scales in the framework of classically scale-invariant theories. We then apply the method to generate the Planck mass and to solve the neutrino mass and dark matter problems through what may be termed the “scale-invariant standard model”. The model is perturbatively renormalisable for energy scales up to the Planck mass.

1 Introduction

It is quite remarkable that almost all the mass of the visible matter in the universe originates from quantum effects that trigger dimensional transmutation in QCD, even though that theory, at the classical level, is strictly invariant under scale transformations when quark masses are neglected. The idea that all the elementary particles, including those constituting dark matter, obtain their masses through the mechanism of dimensional transmutation is therefore very appealing. Indeed, a perturbative mechanism for electroweak symmetry breaking in classically scale-invariant models was presented a long time ago by Coleman and Weinberg [1] (see also [2]).¹ Moreover, classical scale invariance can serve as the symmetry that ensures the stability of the electroweak scale under radiative corrections [3]–[6].

Recently, a number of scale-invariant particle physics models have been proposed [5]–[8]. To be fully realistic, any particle physics model, including those featuring classical scale invariance, must explain neutrino masses, dark matter and the cosmological matter-antimatter asymmetry. This often involves the use of other energy scales, not just that of electroweak symmetry breaking. This presents a particular challenge for classically scale-invariant theories. In fact the above list should then have the Planck mass scale added to it, because a complete scale-invariant theory must include gravity and a mechanism for generating its fundamental scale through the quantum scale anomaly.

In this paper we shall discuss scale-invariant models with stable hierarchically-separated mass scales that are perturbatively renormalisable up to the Planck scale, where the latter feature ensures calculability. We first describe a simple general formalism using hidden sectors that achieves this purpose. We then apply that formalism to generating the Planck scale together with neutrino masses. We shall associate one of the mass scales with the right-handed neutrino Majorana mass scale, thus implementing the see-saw mechanism for light neutrino masses in the context of scale-invariant models. As already noted, the generation of a radiatively stable mass hierarchy in scale-invariant models is not a trivial task.² Because tree-level masses are absent in scale-invariant theories, a hierarchy of mass scales can only be generated through a hierarchy of dimensionless coupling constants. Such a hierarchy will be technically natural — stable under quantum corrections — if sectors which contain the different mass scales decouple from each other in the limit where the relevant coupling constants vanish [6]. Any dark matter sector is “hidden” by definition. We shall demonstrate that our framework can incorporate mirror dark matter, which actually does not require the generation of a new scale.

¹Technicolour models of electroweak symmetry breaking are examples of QCD-like non-perturbative models realising the dimensional transmutation mechanism.

²In [4] a see-saw model for neutrino masses within the scale-invariant standard model is discussed. However, the hierarchy between the electroweak and see-saw scales in [4] is actually based on large-log radiative corrections, thus undermining the whole perturbative approach. Another attempt was made in [7], but the hierarchy there was due to a small Yukawa coupling constant rather than because a genuinely new mass scale was generated.

2 Generating a stable hierarchy of scales

Consider a hidden (high mass scale) sector consisting of a set of scalar fields (S_1, S_2, \dots) which are singlets under the standard model (SM) gauge group. In the limit where these scalars decouple from the visible sector involving the SM fields, and, in particular, the standard electroweak Higgs doublet field ϕ , the scalar potential separates:

$$V(\phi, S_1, S_2, \dots) = V(\phi) + V(S_1, S_2, \dots). \quad (1)$$

In this limit $V(\phi)$ is simply the Coleman-Weinberg potential, and given the heavy top quark, spontaneous symmetry breaking does not arise. Thus we have a massless Higgs particle in this limit. However, the $V(S_1, S_2, \dots)$ part can induce spontaneous symmetry breaking, leading to $\langle S_j \rangle \neq 0$. Now, if we allow the hidden sector to couple to the ordinary matter sector via

$$V = \sum_i \lambda_x^i \phi^\dagger \phi S_i^2, \quad (2)$$

then the symmetry breaking will be communicated to the electroweak sector. Indeed, for some negative $\sum_i \lambda_x^i \langle S_i^2 \rangle$ the interactions in (2) trigger nonzero vacuum expectation value (VEV) $\langle \phi^2 \rangle = -\frac{1}{\lambda_\phi} \sum_i \lambda_x^i \langle S_i^2 \rangle$, where λ_ϕ is the electroweak Higgs self-interaction coupling. The hierarchy of scales $\langle \phi \rangle / \langle S_j \rangle$ is thus controlled by the adjustable parameters λ_x^i . The quantum corrections to the light mass scale $\langle \phi \rangle$ due to the heavy mass scales $\langle S_j \rangle$ are also entirely controlled by the same coupling constants λ_x^i . In the limit $\lambda_x^i \rightarrow 0$ the heavy and light sectors decouple from each other, and $\langle \phi \rangle / \langle S_j \rangle \rightarrow 0$. Therefore, no radiative correction can significantly disturb the light mass scale $\langle \phi \rangle$, and, hence, we have a technically natural solution to the hierarchy problem. Technically natural hierarchies can be similarly generated within the hidden sector of S -fields as well.

The simplest case of a hidden sector consisting of just one real scalar S was discussed earlier [6]. In that case, λ_x induces symmetry breaking in both the SM and hidden sectors. A drawback of this scenario is that the Higgs mass must be relatively large, $M_H^2 \gtrsim \sqrt{12} M_t^2$, which means that the model does not remain perturbative up to the Planck scale, and also is not consistent with the constraint from precision electroweak data. This motivates consideration of the next simplest model consisting of a Higgs doublet and two real scalars, S_1 and S_2 , which we consider below. We shall show that such a model is perturbative up to the Planck scale and consistent with constraints from precision electroweak data. The model can also naturally incorporate neutrino masses via the see-saw mechanism. Dark matter can be introduced by extending the hidden sector to include a mirror copy of all the known particles, which we also discuss.

3 The two-scalar-singlet model

Let us start by working in the decoupled limit, and just consider the hidden sector tree-level potential $V_0(S_1, S_2)$,

$$V_0(S_1, S_2) = \frac{\lambda_1}{4} S_1^4 + \frac{\lambda_2}{4} S_2^4 + \frac{\lambda_3}{2} S_1^2 S_2^2, \quad (3)$$

where for simplicity we impose invariance under $S_1 \rightarrow -S_1$. All mass terms and coupling terms other than the quartic are zero at the classical level because of the imposed scaling symmetry.

Parameterising the fields through

$$S_1 = r \cos \omega, \quad S_2 = r \sin \omega, \quad (4)$$

the potential (3) is then rewritten as

$$V_0(r, \omega) = \frac{r^4}{4} (\lambda_1 \cos^4 \omega + \lambda_2 \sin^4 \omega + 2\lambda_3 \sin^2 \omega \cos^2 \omega). \quad (5)$$

This tree-level potential is quantally corrected as per the Coleman-Weinberg analysis. We shall work in the parameter regime where the one-loop perturbative correction $\delta V_{1\text{-loop}}$ is accurate. Ideally, one would like to directly minimise the corrected potential $V \simeq V_0 + \delta V_{1\text{-loop}}$, but this is impossible to do analytically. We instead follow the approximate procedure introduced by Gildener and Weinberg [2] which is valid in our weakly-coupled theory.

The procedure requires us to at first partially ignore the radiative corrections and to minimise the tree-level potential (5), but with the recognition that in the quantum theory the parameters $\lambda_{1,2,3}$ become running coupling constants, depending on renormalisation scale μ . The tree-level potential has flat radial directions, which means we begin by taking $\langle r \rangle$ to be arbitrary (but nonzero). There are two possibilities, depending on the sign of λ_3 . For the case $\lambda_3 > 0$, the possible symmetry breaking patterns are,

$$\begin{aligned} \sin \omega = 0 \text{ with } \lambda_1(\Lambda) = 0 &\Rightarrow \langle S_1 \rangle = \langle r \rangle = v, \langle S_2 \rangle = 0 \text{ or} \\ \cos \omega = 0 \text{ with } \lambda_2(\Lambda) = 0 &\Rightarrow \langle S_2 \rangle = \langle r \rangle = v, \langle S_1 \rangle = 0. \end{aligned} \quad (6)$$

Taking the first solution for definiteness, the scale Λ is the renormalisation point where λ_1 vanishes. The dimensionless parameter λ_1 transmutes into the scale Λ in the quantised theory. This is a manifestation of the scale anomaly of quantum field theory which generates dimensionful quantities such as masses despite the classical scale invariance.

For the case $\lambda_3 < 0$, both S_1 and S_2 gain VEVs,

$$\langle S_1 \rangle = \langle r \rangle \left(\frac{1}{1 + \epsilon} \right)^{1/2} \equiv v, \quad \langle S_2 \rangle = v\epsilon^{1/2}, \quad (7)$$

where

$$\langle \tan^2 \omega \rangle \equiv \epsilon = \sqrt{\frac{\lambda_1(\Lambda)}{\lambda_2(\Lambda)}}, \quad (8)$$

with

$$\lambda_3(\Lambda) + \sqrt{\lambda_1(\Lambda)\lambda_2(\Lambda)} = 0 \quad (9)$$

imposed. As in the previous case, this relation between the Higgs potential parameters serves to define the renormalisation point, and a dimensionless parameter is transmuted into the scale Λ .

The hierarchy between the VEVs of S_1 and S_2 is determined through the parameter ϵ .³ We can thus immediately apply the above to the hierarchy between the Planck mass and the see-saw scale, so that

$$\epsilon \sim (M_{\text{see-saw}}/M_P)^2, \quad (10)$$

where S_1 is required to couple to the gravitational scalar curvature via $\mathcal{L} \sim \sqrt{-g}S_1^2 R$, and its VEV therefore generates the Newton constant [6]. The smaller VEV of S_2 generates masses for right-handed neutrinos through the Yukawa couplings,⁴

$$\mathcal{L}_{\text{Majorana}} = \lambda_M^i \bar{\nu}_R^i (\nu_R^i)^c S_2 + \text{H.c.} \quad (11)$$

We next calculate the tree-level Higgs masses. We first define the shifted fields $S_1 = \langle S_1 \rangle + S'_1$, $S_2 = \langle S_2 \rangle + S'_2$, and substitute them into the potential, Eq. (3). Of the two physical scalars only one (call it $S = \sin \omega S_1 - \cos \omega S_2$) gains mass at tree-level,

$$\begin{aligned} m_S^2 &= \lambda_3 v^2 & \text{when } \lambda_3 > 0 \\ m_S^2 &= 2(\lambda_1 - \lambda_3)v^2 & \text{when } \lambda_3 < 0. \end{aligned} \quad (12)$$

The other one (call it $s = \cos \omega S_1 + \sin \omega S_2$) remains massless due to a flat direction in the Higgs potential. It is the pseudo-Goldstone boson (PGB) of anomalously-broken scale invariance.

The PGB gains mass from quantal corrections. The 1-loop correction along the flat direction in V_0 is [2]

$$\delta V_{1\text{-loop}} = Ar^4 + Br^4 \log \left(\frac{r^2}{\Lambda^2} \right), \quad (13)$$

where

$$\begin{aligned} A &= \frac{1}{64\pi^2 \langle r \rangle^4} \left[3\text{Tr} \left(M_V^4 \log \left(\frac{M_V^2}{\langle r \rangle^2} \right) \right) \right. \\ &\quad \left. + \text{Tr} \left(M_S^4 \log \left(\frac{M_S^2}{\langle r \rangle^2} \right) \right) - 4\text{Tr} \left(M_F^4 \log \left(\frac{M_F^2}{\langle r \rangle^2} \right) \right) \right], \end{aligned} \quad (14)$$

and

$$B = \frac{1}{64\pi^2 \langle r \rangle^4} [3\text{Tr} M_V^4 + \text{Tr} M_S^4 - 4\text{Tr} M_F^4]. \quad (15)$$

The traces go over all internal degrees of freedom, with $M_{V,S,F}$ being the tree-level masses respectively for vectors, real scalars and Dirac fermions evaluated for the given VEV pattern. In this simple case of just two scalars, the only masses we need to consider are the scalar S and the Majorana fermions ν_R (which means the 4 multiplying M_F^4 in Eq. (15) becomes a 2).

³Note that the hierarchy between the VEVs of S_1 and S_2 is stable under radiative corrections. The small parameter λ_1 receives a 1-loop correction that is proportional to λ_3^2 , which in turn is equal to $\lambda_1 \lambda_2$ when evaluated at the scale Λ . Thus the correction to λ_1 is under control provided λ_2 is not too large, a condition we need in any case to make our weak-coupling analysis valid.

⁴We are not interested in the flavour structure of neutrino mass matrices here, so consider flavour-diagonal couplings for simplicity. Note that the discrete $S_1 \rightarrow -S_1$ symmetry prevents couplings of S_1 to right-handed neutrinos, but in its absence S_1 would also contribute to right-handed neutrino masses.

The extremal condition $\frac{\partial \delta V_{1\text{-loop}}}{\partial r}|_{r=\langle r \rangle} = 0$ tells us that

$$\log \left(\frac{\langle r \rangle}{\Lambda} \right) = -\frac{1}{4} - \frac{A}{2B}. \quad (16)$$

The PGB mass is then [2]:

$$\begin{aligned} m_s^2 &= \frac{\partial^2 \delta V_{1\text{-loop}}}{\partial r^2}|_{r=\langle r \rangle} = 8B\langle r \rangle^2 \\ &= \frac{1}{8\pi^2\langle r \rangle^2} [3\text{Tr} M_V^4 + \text{Tr} M_S^4 - 4\text{Tr} M_F^4] . \end{aligned} \quad (17)$$

Applying this general formula to our theory we obtain,

$$m_s^2 \simeq \frac{1}{8\pi^2\langle r \rangle^2} \left[m_S^4 - 2 \sum m_{\nu_R^i}^4 \right] \quad (18)$$

Now, let us make the model realistic by introducing the Higgs doublet ϕ into the picture. Consider the Higgs potential terms,

$$V_0(\phi, S_1, S_2) = \frac{\lambda_\phi}{2}(\phi^\dagger \phi)^2 + \frac{1}{2}\phi^\dagger \phi(\lambda_{x1}S_1^2 + \lambda_{x2}S_2^2) . \quad (19)$$

If $\lambda_x < 0$ then $\langle \phi \rangle \neq 0$ is induced, with

$$\langle \phi \rangle^2 = -\frac{\lambda_{x1}}{2\lambda_\phi} \langle S_1 \rangle^2 - \frac{\lambda_{x2}}{2\lambda_\phi} \langle S_2 \rangle^2 , \quad (20)$$

where $\langle \phi \rangle = v_{\text{EW}} \approx 174 \text{ GeV}$.⁵ In the decoupling limit $\lambda_{xi} \rightarrow 0$ and $\langle \phi \rangle / \langle S_i \rangle \rightarrow 0$. The standard see-saw mechanism requires the introduction of couplings between left-handed and massive right-handed neutrinos,

$$\mathcal{L}_{\text{Dirac}} = \lambda_D^i \bar{f}_L^i \phi \nu_R^i + \text{H.c.} \quad (21)$$

These couplings together with the couplings in (11) generate quantal corrections to λ_{x2} , so that technically the decoupling limit corresponds to $\lambda_{xi} \rightarrow 0$ and $\lambda_M^i \lambda_D^i \rightarrow 0$.

One can obtain a naturalness constraint on the couplings λ_M^i, λ_D^i by demanding that there be no fine-tuned cancellation between the tree level and the 1-loop contribution to λ_{x2} . This condition implies,

$$\frac{(\lambda_M^i \lambda_D^i)^2}{16\pi^2} \lesssim \lambda_\phi \frac{v_{\text{EW}}^2}{\langle S_2 \rangle^2}. \quad (22)$$

⁵The solutions in Eqs. (6) and (7) are modified due to the additional terms in Eq. (19). However, because of the hierarchy $\langle \phi \rangle / \langle S_i \rangle \ll 1$, these modifications are insignificant and we have neglected them.

This condition can be rewritten in terms of the physical masses:

$$M^i \lesssim \left(\frac{16\pi^2 v_{\text{EW}}^4}{m_\nu^i} \right)^{1/3}. \quad (23)$$

The most stringent constraint on M^i comes from the most massive neutrino, which from the atmospheric neutrino anomaly suggests that the lightest M^i satisfies $M^i \lesssim 10^7$ GeV. A similar bound has been obtained in [9].

Note that all the couplings in the Higgs potential can be $\lll \mathcal{O}(1)$, except for λ_ϕ . This means that contributions to the renormalization group beta-functions from interactions with hierarchically small coupling constants are negligible compared to other relevant contributions steaming from the Standard Model interactions. Namely, the evolution of the electroweak Higgs self-interaction coupling will be governed by $\beta_{\lambda_\phi} \simeq \beta_{\lambda_\phi}^{\text{SM}}$, where $\beta_{\lambda_\phi}^{\text{SM}}$ is the Standard Model beta-function. This evolution is largely unaffected by the presence of the heavy hidden sectors simply because the interactions between the sectors are assumed to be extremely weak. Therefore, the theory is manifestly perturbative up to the Planck scale, so long as the Higgs mass is relatively light. Moreover, boundedness of the potential requires positivity of λ_ϕ for all energy scales up to the Planck mass, which gives a lower bound on the Higgs boson mass. The range of the allowed Higgs boson masses is established through the solutions of the corresponding renormalisation group equations. In the regime of dominant λ_ϕ coupling the results are similar to those obtained within the minimal standard model [10]:

$$129 \text{ GeV} \lesssim m_h \lesssim 175 \text{ GeV}. \quad (24)$$

4 Incorporating mirror dark matter

The mirror extension of the standard model [11] is one of the best motivated models of a hidden sector. The SM gauge group G_{SM} is extended to $G_{SM} \times G'_{SM}$, where the second factor is isomorphic to the first but independent of it. Standard particles are singlets under G'_{SM} , while their mirror partners are singlets under G_{SM} . A discrete parity symmetry that interchanges the two sectors is imposed. The fact that it is a parity symmetry means that a left-handed standard fermion is partnered by a right-handed mirror fermion, and so on. When the see-saw model is extended by a mirror sector, the usual right-handed gauge-singlet neutrino ν_R is partnered by a left-handed gauge-singlet neutrino ν'_L , where primes denote mirror partners. They interchange under mirror parity. The standard Higgs doublet is partnered by a mirror Higgs doublet.

Mirror models naturally incorporate dark matter candidates, because if a given ordinary atom is stable then so is its mirror partner. It is remarkable that, unlike many popular dark matter models, the simplest model with unbroken mirror symmetry is capable of explaining [12, 13] both the DAMA [14] and CoGeNT [15] data and also the null results of the other experiments. In [5] we considered the minimal scale-invariant mirror matter model and demonstrated that radiative electroweak symmetry breaking is possible due to the presence of extra bosonic degrees of freedom, in particular the mirror Higgs boson. However, the minimal model also requires spontaneous mirror

symmetry breaking and is plagued with a Landau pole problem at energies much below the Planck mass. In addition, neutrinos are massless in the minimal model.

All these problems can be solved within a simple extension of the minimal model of Ref. [5]. It is obtained as a variation of the model considered in the previous section, where now the hidden sector is extended to a full mirror sector. The scalars S_1 and S_2 are assigned as singlets of mirror parity, that is $S_i \rightarrow S_i$. The field S_2 couples to the ν'_L 's as well as to the ν_R 's,

$$\mathcal{L}_{\text{Majorana}} = \lambda_M^i [\bar{\nu}_R^i (\nu_R^i)^c + \bar{\nu}_L'^i (\nu_L'^i)^c] S_2 + \text{H.c.} \quad (25)$$

thus generating the see-saw scale in both sectors. The scalar potential comprises of (3) and

$$\begin{aligned} V_0(\phi, \phi', S_i) &= \frac{\lambda_\phi}{2} [(\phi^\dagger \phi)^2 + (\phi'^\dagger \phi')^2] + \frac{\lambda_{\phi\phi}}{2} (\phi^\dagger \phi)(\phi'^\dagger \phi') \\ &+ \frac{1}{2} (\phi^\dagger \phi + \phi'^\dagger \phi') (\lambda_{x1} S_1^2 + \lambda_{x2} S_2^2). \end{aligned} \quad (26)$$

Besides the VEVs that spontaneously generate the see-saw and Planck scales (7)-(9)⁶, one obtains, for a finite region of parameter space, mirror symmetry preserving VEVs that break electroweak and mirror-electroweak gauge invariance:

$$\langle \phi \rangle = \langle \phi' \rangle \equiv v_{\text{EW}}^2 = -\frac{\lambda_{x1} + \epsilon \lambda_{x2}}{2\lambda_\phi + \lambda_{\phi\phi}} v^2. \quad (27)$$

Again, the hierarchy $v_{\text{EW}}^2/v^2 \sim M_Z^2/M_P^2 \lll 1$ is determined through the hierarchy of dimensionless couplings, $-\frac{\lambda_{x1} + \epsilon \lambda_{x2}}{2\lambda_\phi + \lambda_{\phi\phi}} \lll 1$. This hierarchy is technically natural because of the decoupling limit $\lambda_{xi} \rightarrow 0$.

5 Conclusion

The mass carried by the visible and dark matter in the universe may have intrinsically quantum origin as implemented in classically scale-invariant theories. We have discussed a class of simple scale-invariant models which incorporate a stable hierarchy of mass scales, in particular the hierarchy between the Planck, neutrino see-saw and the electroweak scales. We have found that the mass of the lightest right-handed neutrino can be as large as $\sim 10^7$ GeV, without disturbing the electroweak scale. This naturalness bound on the lightest right-handed neutrino suggests that the simplest leptogenesis [16] scenario is disfavoured due to the insufficient CP-violation [17]. However, there are other possibilities such as the resonant leptogenesis scenario (see e.g. [18]) with nearly degenerate right-handed neutrinos, which remain a viable option. Within the simplest model, perturbativity and stability requirements bound the Higgs boson mass to the range given in Eq. (24). Within the same model one can also simultaneously generate the Planck mass. By

⁶Note, that while the mass of heavy scalar S is still given by Eq. (12), the mass of lighter scalar s is modified due to the extra contribution from right-handed mirror neutrinos. Thus, instead of (18), we have now $m_s^2 \simeq \frac{1}{8\pi^2 \langle r \rangle^2} [m_S^4 - 4 \sum m_{\nu_R^i}^4]$.

extending the hidden sector to a full mirror sector, we have obtained a viable scale-invariant mirror model with unbroken mirror symmetry. In this way one can incorporate dark matter which is in remarkable agreement [12], [13] with recent experimental results from the DAMA [14] and CoGeNT [15] collaborations.

Acknowledgements

This work was supported in part by the Australian Research Council.

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